

# On Electroweak Baryogenesis in Gauge Mediated Models with Messenger-Matter Mixing

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## Abstract

We consider feasibility of electroweak baryogenesis in gauge mediated supersymmetry breaking models with messenger-matter mixing. We present a class of models where electroweak baryogenesis produces sufficient amount of baryon asymmetry. The main features of these models are (i) large mixing between messengers and right stop and (ii) fairly narrow region of viable  $\tan\beta$ .

## 1 Introduction

Among various puzzles of Nature the problem of baryogenesis is one of the most intriguing. Since pioneering works [1], necessary conditions for the generation of the baryon asymmetry are well-known. These are baryon number violation, C- and CP-nonconservation and departure from thermal equilibrium at a certain stage of the evolution of the Universe. Search for reliable mechanisms of baryogenesis is limited to theories where these conditions are met.

One of the most appealing mechanisms is electroweak baryogenesis [2, 3]. It was proposed that this mechanism works during the electroweak phase transition. Unlike other mechanisms, electroweak baryogenesis was thought to be inherent in the Standard Model (SM) and seemed to require no additional fields. However, the baryon asymmetry tends to be washed out by anomalous electroweak processes if the latter come into thermal equilibrium after the phase transition completes. To prevent this wash-out, the electroweak phase transition has to be of the (strong enough) first order; this imposes a constraint on the expectation value of Higgs field  $v(T)$  (with  $v(0) \approx 245$  GeV) at the critical temperature  $T_c$  [4],

$$\frac{v(T_c)}{T_c} \gtrsim 1. \quad (1)$$

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In the case of SM this requirement implies the bound on the Higgs boson mass [3],

$$m_h < 40 \text{ GeV} , \quad (2)$$

which is inconsistent with present experimental limits [5]. Hence electroweak baryogenesis fails in the Standard Model.<sup>1</sup>

Similar situation takes place in the Minimal Supersymmetric Standard Model (MSSM). Although the corresponding bounds on the parameters of MSSM are significantly weaker, almost entire parameter space consistent with successful electroweak baryogenesis is excluded by existing experimental data. Still, electroweak baryogenesis works in MSSM with light Higgs, light  $\tilde{t}_R$ , heavy  $\tilde{t}_L$ , small  $\tilde{t}_R - \tilde{t}_L$  mixing, and charginos with fairly degenerate masses [7] (for brief reviews, see Refs. [8, 9]). This spectrum was discussed in the framework of models with gravity mediation of supersymmetry breaking [7], and it was observed that the price for this “light stop” solution is non-universal boundary conditions for soft terms at GUT scale.

Recently, the phenomenology of Gauge Mediated Supersymmetry Breaking (GMSB) models attracted considerable attention (for a review, see Ref. [10]). In these models, supersymmetry breaking occurs in a separate sector. Supersymmetry breaking terms in the visible sector are generated due to special fields (messengers) charged under SM gauge group. The soft masses of superpartners are determined by their quantum numbers and are proportional to the corresponding gauge coupling constants. Consequently, all squarks are very heavy and electroweak baryogenesis seems not to work. So, it has been argued that the favorite mechanism to produce the baryon asymmetry in GMSB models is the Affleck–Dine baryogenesis [11].

This letter addresses the possibility of electroweak baryogenesis in GMSB with messenger-matter mixing. This mixing is natural since messengers carry the same quantum numbers as ordinary fields in MSSM [12]. The constraints on mixing parameters imposed by the absence of lepton flavor violating processes and FCNC are not particularly strong [13, 14]. Moreover, it was pointed out that some scalar masses may be significantly reduced at large values of mixing terms without any contradiction to experimental limits [14]. We will see that the small mass of stop required for successful electroweak baryogenesis may be explained in this way. We will show also that all other conditions on the spectrum may be satisfied.

## 2 Light stop window for electroweak baryogenesis

In any model of baryogenesis two main questions arise: *i*) What is the mechanism of the generation of the baryon asymmetry? *ii*) What protects the baryon asymmetry from being washed out by sphaleron processes?

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<sup>1</sup>There are recent attempts (see, e.g., Ref. [6]) to revive electroweak baryogenesis in SM by invoking the dynamics of preheating.

In the framework of “light stop” solution, the answers to these questions are as follows.

*i)* The relevant source of CP-violation is the phase  $\phi_{CP}$  of  $\mu$ -term in Higgs superpotential. Baryogenesis is fueled by CP-asymmetry in chargino flux through expanding bubble wall. The realistic amount of baryon asymmetry  $n_B/s$  is produced by sphaleron processes provided that chargino and neutralino masses are not much larger than the critical temperature  $T_c \sim 100$  GeV. The resulting asymmetry is determined by the ratios of masses to the critical temperature, and degenerate  $\mu$  and wino mass  $M_2$  are favorable. Though the general tendencies are clear, the actual calculated values of  $n_B/s$  depend on the approximations made. In what follows we make use of the constraints on the values of CP-phase and the level of degeneracy presented in Ref. [15]. The smallest allowed value of  $\phi_{CP}$  is  $|\phi_{CP}| = 10^{-4}$ ; at this value one requires  $\mu = M_2$ . In the opposite case of large CP-violating phase,  $|\sin \phi_{CP}| = 1$ , the allowed region of masses is  $\mu = \epsilon M_2$ ,  $0.4 \lesssim \epsilon \lesssim 2.5$  at  $\mu, M_2 \simeq T_c$  and  $0.5 \lesssim \epsilon \lesssim 2$  at  $\mu, M_2 \simeq 3T_c$ . The favorable region of the mass of CP-odd Higgs boson is  $m_A \lesssim 300$  GeV, otherwise  $n_B/s$  is suppressed by  $m_A^{-2}$  [16, 8]. For heavier CP-odd Higgs, the values of  $\mu$  and  $M_2$  are to be more degenerate in order that the baryon asymmetry at given  $\phi_{CP}$  be the same as at small  $m_A$ .

*ii)* As mentioned above, electroweak phase transition has to be of the first order, so that the constraint (1) is satisfied. Light stop helps in strengthening the phase transition. The largest baryon asymmetry is obtained in the context of MSSM with right stop plasma mass vanishing at the critical temperature, when the condition (1) becomes [17]

$$1 \lesssim \frac{v(T_c)}{T_c} = \left( \frac{v(T_c)}{T_c} \right)_{SM} + \frac{2m_t^3 \left( 1 - \frac{\tilde{A}_t^2}{m_Q^2} \right)^{3/2}}{\pi v m_h^2}, \quad (3)$$

where  $\tilde{A}_t = |A_t - \mu \cot \beta|$ , and  $A_t$  is the stop trilinear soft term,  $m_Q$  is the left stop mass and  $m_t$  is the on-shell running top quark mass in the  $\overline{MS}$  scheme. Hereafter  $m_h$  denotes the mass of the lightest Higgs boson. We consider the case  $m_A \gg T_c$ , which is relevant to GMSB models; in this case CP-odd Higgs does not affect the electroweak phase transition. The first term on the right hand side of Eq. (3) is the ordinary Standard Model contribution,

$$\left( \frac{v(T)}{T} \right)_{SM} \simeq \left( \frac{40 \text{ GeV}}{m_h} \right)^2.$$

The second term in Eq. (3) is the contribution from light right stop. Current limits on the lightest Higgs mass [5],  $m_h \gtrsim 100$  GeV, combined with the inequality (3) impose a constraint on left-right mixing

$$\tilde{A}_t/m_Q \lesssim 0.5. \quad (4)$$

At larger  $m_h$  left-right mixing has to be smaller. At  $\tilde{A}_t = 0$  Eq. (3) gives the upper bound on the lightest Higgs boson mass in the theory with successful electroweak baryogenesis,  $m_h \simeq 115$  GeV [9]. It is worth noting that the result (3) has been obtained by making use of improved one-loop effective potential. Higher order corrections make the phase transition slightly stronger (for a brief review and references, see Ref. [9]) and will be neglected in what follows.

Zero right stop plasma mass at the critical temperature implies

$$m_{\tilde{t}_R}^{2(eff)} = m_U^2 + \Pi_R(T_c) \approx 0, \quad (5)$$

where  $m_U$  is the low energy value of the right stop soft mass term and  $\Pi_R(T_c)$  is the finite temperature contribution to the effective squared mass which is of order  $T_c^2$  [7]. Hence one needs negative  $m_U^2 \sim -(100 \text{ GeV})^2$ , which in principle may result in the existence of charge- and color-breaking (local) vacuum. A conservative requirement is that the physical vacuum has lower energy than the color-breaking minimum. The latter condition at, as an example,  $\tilde{A}_t = 0$  yields approximately [7]

$$|m_U| < m_{crit} = \left( \frac{m_h^2 v^2 \alpha_3 \pi}{3} \right)^{1/4}, \quad (6)$$

that is  $|m_U| \lesssim 95$  GeV for  $m_h = 100$  GeV. One can relax this constraint by considering the physical vacuum as a metastable but long-living minimum. However, the inequality

$$m_U^2 + \Pi_R(T_c) > 0 \quad (7)$$

has to be satisfied in any case, otherwise the Universe would be driven to a charge- and color-breaking minimum at  $T > T_c$  (for a discussion see Ref. [7]).

Let us collect the requirements which are imposed on the theory with “light stop” solution:

1. Right-left mixing in stop sector is small,  $\tilde{A}_t < 0.5m_Q$  at  $m_h > 100$  GeV, and the heavier the lightest Higgs boson, the smaller the mixing.
2. At  $m_h > 115$  GeV there is no window for electroweak baryogenesis. Consequently, the mass of the lightest Higgs boson belongs to the interval  $100 \text{ GeV} < m_h < 115$  GeV, where the lower bound comes from experiment.
3. The favorable interval of the mass of CP-odd Higgs boson is  $150 \text{ GeV} \lesssim m_A \lesssim 300$  GeV. At smaller  $m_A$  the phase transition is weaker, while at larger  $m_A$  the baryon asymmetry is suppressed by  $m_A^{-2}$ .
4. The soft mass squared of right stop is negative and of order  $-T_c^2$ . Inequalities (6), (7) (or similar ones at non-zero  $\tilde{A}_t$ ) have to be satisfied.

### 3 Gauge Mediated Models with messenger-matter mixing

Let us recall the spectrum of superpartners in GMSB models [10]. In simple versions, messengers belong to complete vector-like GUT (e.g.,  $SU(5)$ ) multiplets. The induced soft terms depend crucially on  $n = n_5 + 3n_{10}$ , with  $n_5$  and  $n_{10}$  being the numbers of  $5 + \bar{5}$  and  $10 + \bar{10}$  messenger generations. Indeed, the spectrum of superpartners in these models at the messenger scale  $M_m$  is

$$M_i = nc_i \frac{\alpha_i}{4\pi} \Lambda f_1\left(\frac{\Lambda}{M_m}\right), \quad (8)$$

$$m^2 = 2n\Lambda^2 \sum_{i=1}^3 C_i c_i \left(\frac{\alpha_i}{4\pi}\right)^2 f_2\left(\frac{\Lambda}{M_m}\right), \quad (9)$$

where  $M_i$  denote gaugino masses and  $m$  are soft masses of the scalar superpartners of fermions of the Standard Model. Here  $\alpha_1 = \alpha/\cos^2\theta_W$ ,  $\alpha_2 = \alpha/\sin^2\theta_W$ ,  $\alpha_3$  are gauge coupling constants of electroweak and strong interactions;  $c_1 = 5/3$ ,  $c_2 = c_3 = 1$ ;  $C_3 = 4/3$  for color triplets (zero for singlets),  $C_2 = 3/4$  for weak doublets (zero for singlets),  $C_1 = \left(\frac{Y}{2}\right)^2$ , where  $Y = 2(Q_{em} - T_3)$  is the weak hypercharge, and  $\Lambda < M_m$  is a dimensional parameter related to the scale of supersymmetry breaking. The functions  $f_1$  and  $f_2$  weakly depend on their argument and are close to 1 in the most part of the domain of definition. We consider soft mixing term in the Higgs sector  $B_\mu$  (or  $\tan\beta$ ) as a free parameter. The value of  $\mu$  is determined by the relation

$$\mu^2 = -\frac{1}{2}M_Z^2 + \frac{m_{h_D}^2 - m_{h_U}^2 \tan^2\beta}{\tan^2\beta - 1}, \quad (10)$$

where  $m_{h_U}$  and  $m_{h_D}$  are soft masses of up- and down-Higgs bosons. Note that gaugino masses grow as  $n$ , whereas scalar ones grow as  $\sqrt{n}$ . It is this behavior that enables one to obtain the degeneracy  $\mu \sim M_2$  favorable for successful generation of baryon asymmetry.

Electroweak symmetry breaking leads to additional contributions to the soft mass spectrum. D-terms drive the masses up while mixing between left and right scalars increases the splitting between their masses. Taking into account these corrections and the renormalization group evolution from  $M_m$  to  $M_{SUSY}$  one calculates the low energy spectrum. In fact, reasonable estimates for the low energy soft masses are obtained by plugging  $\alpha_i = \alpha_i(M_{SUSY})$  into Eqs. (8) and (9). As concerns the trilinear soft terms, their values at  $M_m$  are additionally suppressed by coupling constants in comparison with Eqs. (8), (9) and become significant at low energies only due to the renormalization group evolution. The largest low energy trilinear coupling is  $A_t$ , which is numerically  $A_t \simeq M_2$ .

We consider messengers which are odd under R-parity. Then the components of the fundamental messengers have the same quantum numbers as left leptons and right down-quarks, while components of antisymmetric messengers have quantum numbers of right leptons, left quarks and right up-quarks. Therefore, symmetries allow for mixing between messengers and matter fields [12]. In the case of one fundamental messenger multiplet, the mixing terms in the superpotential are

$$\mathcal{W}_{mm}^{(5)} = H_D L_m Y_i^{(5)} E_i + H_D D_m X_i^{(5)} Q_i , \quad (11)$$

where  $i = 1, 2, 3$  counts matter generations and subscript  $m$  refers to messenger fields. We use the standard notations for MSSM superfields ( $E_i$  are right leptons and  $Q_i$  are left up-quarks),  $Y_i^{(5)}$  and  $X_i^{(5)}$  are mixing parameters,  $L_m$  and  $D_m$  are messenger superfields with quantum numbers of left leptons and right down-quarks, respectively.

When messengers are integrated out, mass matrices of scalar fields acquire negative contributions,

$$\delta m_{ij}^2 \sim -10^{-2} \Lambda^2 Y_i^* Y_j ,$$

that emerge from one-loop diagrams with messenger fields running in loops [12] (here  $Y_i$  generically stands for either  $Y_i^{(5)}$  or  $X_i^{(5)}$ ). In principle, these terms give rise to flavor violating processes ( $\mu \rightarrow e\gamma$ ,  $b \rightarrow s\gamma$ , etc.); the corresponding limits on  $Y_i$  are derived in Ref. [14]. It is important for our discussion that the mixing terms also reduce some scalar masses. Indeed, let us consider colorless sector and neglect left-right mixing. Then the eigenvalues of right slepton mass matrix are

$$\{m_R^2, \quad m_R^2, \quad m_R^2 - \sum_{i=1}^3 \delta m_{ii}^2\} ,$$

where  $m_R$  is the soft mass of right slepton. The squark masses are modified in a similar way. This hints towards the possibility to have the right stop soft mass term  $m_{t_R}^2$  negative, as required by the “light stop” scenario of electroweak baryogenesis. However, fundamental messengers are not suitable for this purpose, as they do not mix with right stop. Hence, messenger fields belonging to antisymmetric representation are to be involved.

## 4 An example of GMSB model with “light stop” window

As a concrete example of a model where all criteria of “light stop” solution are satisfied, we consider GMSB model with  $n_{10}$  antisymmetric messengers and non-zero coupling  $Y_3$  just between right stop  $U_3$  and corresponding messengers  $Q_m$ ,

$$\mathcal{W}_{mm}^{(10)} = Y_3 H_U U_3 Q_m . \quad (12)$$

When messengers are integrated out, no additional mixing appears in squark matrix, so there is no problem with FCNC<sup>2</sup>. As a result of mixing (12), the right stop mass squared,  $m_{t_R}^2$ , and up-Higgs mass squared,  $m_{h_U}^2$ , get negative contributions [14],

$$\delta m_{t_R}^2 \approx -n \frac{\Lambda^2}{8\pi^2} |Y_3|^2 f_3\left(\frac{\Lambda}{M_m}\right), \quad \delta m_{h_U}^2 \approx -n \frac{3\Lambda^2}{16\pi^2} |Y_3|^2 f_3\left(\frac{\Lambda}{M_m}\right) = \frac{3}{2} \delta m_{t_R}^2, \quad (13)$$

where  $f_3(\Lambda/M_m) = (1/6)(\Lambda/M_m)^2$  at  $\Lambda/M_m$  not very close to 1. These terms come from one-loop diagrams involving the Yukawa interaction (12), with messengers running in loops.

For given  $\Lambda$  and  $n = 3n_{10}$  it is possible to tune  $Y_3$  and obtain a negative value of  $m_{t_R}^2$  of order of  $-m_{crit}^2$ . In particular, Eq. (5) is obeyed at

$$|Y_3|^2 \simeq 8\pi^2 \frac{m_{t_R}^2 + m_{crit}^2}{3n_{10}\Lambda^2} f_3^{-1}\left(\frac{\Lambda}{M_m}\right). \quad (14)$$

where  $m_{t_R}^2$  as given by Eq. (9) is typically much larger than  $|m_{crit}|^2$ . The value of the messenger-stop Yukawa coupling turns out to be fairly large. As an example, at  $\Lambda/M_m = 0.5$  one has  $Y_3 \simeq 4\sqrt{2}\alpha_3 \simeq 0.6$ . In this range of  $Y_3$ , Eq. (6) is also straightforward to fulfill. Hence, the fourth requirement of section 2 may be satisfied by tuning  $Y_3$ .

We turn to the other three requirements. The lower bound on the lightest Higgs boson mass is experimental and has to be satisfied regardless of baryogenesis. This is not the only constraint on the parameter space coming from collider experiments: as there is no evidence for light right sleptons, small values  $\Lambda\sqrt{n} < 30$  TeV are forbidden. At the level of our accuracy we approximate the one-loop enhanced mass of the lightest Higgs boson as follows,

$$m_h^2 = M_z^2 \cos^2 2\beta + \frac{3\sqrt{2}}{2\pi^2} G_F m_t^4 \left( \log\left(\frac{m_Q^2}{m_t^2}\right) + \frac{\tilde{A}_t^2}{m_Q^2} \right). \quad (15)$$

Since one of the requirements of the “light stop” solution is small left-right mixing in stop sector (see Eq. (4)), equation (15) implies the lower bound on  $\tan\beta$  depending on the key GMSB parameter,  $\Lambda\sqrt{n}$ , that determines the value of  $m_Q$  through Eq. (9). One obtains  $\tan\beta \gtrsim 3.5$  at  $\Lambda\sqrt{n} = 30$  TeV,  $\tan\beta \gtrsim 2.5$  at  $\Lambda\sqrt{n} = 50$  TeV and  $\tan\beta \gtrsim 1.5$  at  $\Lambda\sqrt{n} = 100$  TeV. Within our approximation, the largest value of  $m_h$  consistent with “light stop” solution,  $m_h < 115$  GeV, is achieved at very large  $\Lambda\sqrt{n}$

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<sup>2</sup>In what follows we will require rather large value of  $Y_3$ . A model with three large mixing terms of the same order of magnitude in right up-squark sector is ruled out by the absence of FCNC. On the other hand, if we assume the hierarchy between messenger-matter couplings similar to the hierarchy between SM fermion Yukawa couplings, we obtain a “natural” model with large messenger–stop mixing and suppressed FCNC.

and hence  $m_Q$ :  $\Lambda\sqrt{n} \gtrsim 400$  TeV,  $m_Q \gtrsim 4$  TeV. The reason is that, as we will see below, successful baryogenesis requires rather small values of  $\tan\beta$ . It is worth noting that similar constraints are imposed on the parameters of SUGRA models with “light stop” solution.

Let us proceed with other requirements. Before discussing concrete cases of small and large CP-violating phase  $\phi_{CP}$ , let us make a few general remarks.

First, mixing (12) drives  $m_{h_U}^2$  deep into the region of negative values,

$$m_{h_U}^2 = m_{h_D}^2 - \frac{3}{8\pi^2} m_Q^2 \left( \log \frac{M_m}{m_Q} + \frac{3}{2} \right) + \frac{3}{2} \delta m_{t_R}^2, \quad (16)$$

that leads to very large  $\mu$  because of Eq. (10). This implies that the CP-phase has to be relatively large, since  $m_A \sim \mu$ , and the baryon asymmetry is suppressed by the mass of heavy CP-odd Higgs boson. Moreover, since an approximate degeneracy  $\mu \simeq M_2$  is required, wino is also heavy. This makes a potential problem, because in GMSB models with complete messenger multiplets one has  $A_t \simeq M_2$ , and there appears large left-right mixing. The cure is to make  $\tilde{A}_t = |A_t - \mu \cot\beta|$  small,  $\tilde{A}_t < 0.5m_Q$ , by choosing suitable sign of  $\mu$  and tuning  $\tan\beta$ .

Second, the degeneracy between  $\mu$  and  $M_2$  will be achieved by a suitable choice of the number of messengers,  $n_{10}$ . Indeed, these two mass parameters scale differently with the number of messenger fields. Namely, from Eq. (10) one finds that  $\mu$  is determined by scalar soft masses, so it grows as  $\sqrt{n}$  (see Eq. (9)), while  $M_2$  grows as  $n$  (see Eq. (8)).

One can estimate viable values of parameters of this model by the following simple chain of reasonings. Squarks are the heaviest scalars in this theory, because the hierarchy of soft terms are governed by corresponding gauge couplings. Since  $\Lambda\sqrt{n} \gtrsim 30$  TeV, the model has heavy strong sector,  $m_Q^2 \gg m_Z^2$ . The “light stop” solution requires  $|m_Q^2 + \delta m_{t_R}^2| \sim M_z^2$ , so  $|\delta m_{t_R}^2|$  is also large, and  $m_{h_U}^2 \approx \frac{3}{2} \delta m_{t_R}^2$  (see Eqs. (16), (13)). Then the relation (10) may be approximated as follows,

$$\mu^2 = \frac{3}{2} m_Q^2 \frac{\tan^2\beta}{\tan^2\beta - 1}. \quad (17)$$

The “light stop” solution requires almost degenerate spectrum,  $\mu = \epsilon M_2$  with  $1/2 < \epsilon < 2$ . In view of the relation  $A_t \simeq M_2$ , Eq. (4) implies

$$\left| \frac{1}{\epsilon} - \frac{1}{\tan\beta} \right| \lesssim \frac{1}{2} \frac{m_Q}{\mu}, \quad (18)$$

Making use of Eq. (17) we obtain from the inequality (18)

$$\left| \frac{1}{\epsilon} - \frac{1}{\tan\beta} \right| \lesssim \frac{1}{\sqrt{6}} \sqrt{1 - \frac{1}{\tan^2\beta}}. \quad (19)$$



This determines the viable values of  $\tan \beta$  for given degeneracy  $\epsilon$ . At  $\epsilon = 1$  ( $M_2 = \mu$ ) one has  $\tan \beta \lesssim 1.5$ , at  $\epsilon = 1.5$  one has  $\tan \beta \lesssim 4$  while at  $\epsilon = 2$  one obtains  $\tan \beta \lesssim 10$ .

By making use of Eq. (17) and Eqs. (8), (9) we find the number of messengers required for successful electroweak baryogenesis at given  $\epsilon$  and  $\tan \beta$ ,

$$n_{10} \simeq \frac{4}{3\epsilon^2} \frac{\alpha_3^2}{\alpha_2^2} \frac{\tan^2 \beta}{\tan^2 \beta - 1}. \quad (20)$$

The estimates are:  $n_{10} \gtrsim 24$  at  $\epsilon = 1$  and  $\tan \beta \lesssim 1.5$ ,  $n_{10} \gtrsim 6$  at  $\epsilon = 1.5$  and  $\tan \beta \lesssim 4$ ,  $n_{10} \gtrsim 3$  at  $\epsilon = 2$  and large  $\tan \beta \lesssim 10$ .

Until now we discussed mostly the bounds coming from the survival of baryon asymmetry after the electroweak phase transition. Additional bounds come from the calculations of the baryon asymmetry along the lines of Ref. [15]. Let us present the results for the cases of small and large  $\phi_{CP}$ . They depend on the value of  $\Lambda$  which determines the mass of the CP-odd Higgs boson  $m_A$ , and hence the suppression factor in the baryon asymmetry.

The smallest CP-phase corresponds to the case of complete degeneracy between  $M_2$  and  $\mu$ , i.e.,  $\mu/M_2 \equiv \epsilon = 1$ . In this case  $\tan \beta \lesssim 1.5$ , so one has to take  $\Lambda\sqrt{n} \gtrsim 100$  TeV in order to obtain  $m_h > 100$  GeV. As discussed above, the number of messengers is large,  $n_{10} \gtrsim 24$ . This results in the large value of Higgsino mass  $\mu \gtrsim 2200$  GeV and large  $m_A \approx \mu/\sin \beta$ , that makes the phase transition stronger, but reduces the amount of baryon asymmetry at given  $\phi_{CP}$ . The realistic value of  $n_B/s$  is obtained at  $\Lambda\sqrt{n} = 100$  TeV with  $|\sin \phi_{CP}| \gtrsim 0.008$ . At  $\Lambda\sqrt{n} = 200$  TeV we find  $\tan \beta \lesssim 1.5$ ,  $n_{10} \gtrsim 24$ ,  $\mu \gtrsim 4400$  GeV and  $|\sin \phi_{CP}| \gtrsim 0.03$ .

In models with large CP-violation phase ( $|\sin \phi_{CP}| = 1$ ), fine tuning between  $M_2$  and  $\mu$  may be somewhat relaxed, and the constraints on  $n_{10}$  and  $\tan \beta$  are, of course, weaker. As an example, let us consider the case  $\mu = 1.5M_2$ . At  $\Lambda\sqrt{n} = 50$  TeV one obtains  $n_{10} \gtrsim 6$ ,  $\mu \gtrsim 800$  GeV and  $\tan \beta \lesssim 4$ , while at  $\Lambda\sqrt{n} = 100$  TeV one gets  $n_{10} \gtrsim 6$ ,  $\mu \gtrsim 1600$  GeV and  $\tan \beta \lesssim 4$ .

Let us note in passing that the case  $\mu < M_2$  is not favorable for the “light stop” solution. Indeed, it follows from Eq. (19) that the viable region of  $\epsilon \equiv \mu/M_2 < 1$  is very narrow and  $\tan \beta$  is smaller than 1.5 at  $\epsilon$  within this interval. Also, small values of  $\Lambda\sqrt{n}$  are disfavored. As an example, at  $\Lambda\sqrt{n} = 30$  TeV one finds  $\mu \gtrsim 450$  GeV, and the absence of right slepton<sup>3</sup> with mass smaller than 45 GeV requires  $\tan \beta \lesssim 3$ . This is in contradiction with the bound coming from the lightest Higgs boson mass,  $\tan \beta \gtrsim 3.5$ .

One concludes that all requirements for the “light stop” solution are satisfied in this model, hence electroweak baryogenesis is capable of generating sufficient amount of baryon asymmetry in GMSB models. The region of viable parameters is fairly

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<sup>3</sup>In the case of large  $\mu$  which is of relevance throughout our discussion the masses squared of right sleptons acquire negative contributions due to electroweak symmetry breaking, which are proportional to  $\mu \tan \beta$  [10].

narrow due to strong restrictions coming from the search for the lightest Higgs boson. Indeed,  $\tan\beta$  in models with “light stop” solution tends to be small, while existing limits on the lightest Higgs boson mass favor high  $\tan\beta$ . At small CP-phase the suitable region stretches along  $\tan\beta \sim 1.5$  and the smallest CP-phase,  $\phi_{CP} \sim 10^{-2}$ , is possible at  $\Lambda\sqrt{n} \simeq 100$  TeV. At  $\phi_{CP} \sim 1$  this region becomes wider and extends to  $\tan\beta \simeq 4$ . The suitable region of parameter space is the largest at the smallest possible  $m_h$ . A drawback of the model is a large number of messenger fields, so the gauge coupling constants become large below the scale of possible Grand Unification.

Electroweak baryogenesis would work better in a model where left squarks are heavy, whereas right squarks are light (or  $\mu$  is relatively small) and  $A_t$  is small. In the framework of GMSB these features appear in models with additional large mixing terms or extended messenger content. Let us briefly discuss these possibilities.

One can introduce additional mixing (e.g., between d-squark and corresponding messengers), which provides large negative contribution to mass squared of down-Higgs. Then larger values of  $\tan\beta$  may become viable. Likewise, the value of  $\mu$  may be smaller, especially at low  $\tan\beta$ , which would extend the region of suitable values of CP-phase.

One can also consider messengers which do not form complete GUT multiplets.<sup>4</sup> Let us denote the effective numbers of messengers generating gluino mass and wino mass as  $n_c$  and  $n_w$ , respectively. Then the relevant quantity will be  $n_w^2/n_c$  instead of  $n_{10}$ . Therefore, it will be possible to reduce significantly the number of messenger fields and relax the problem of unification of gauge coupling constants in a theory with “light stop” window.

There is yet another possibility related to messengers belonging to incomplete GUT multiplets. We outline it by making use of the simplest example of only lepton-like messengers without any messenger-matter mixing. In this model, right stop acquires negative mass squared due to renormalization group evolution, in analogy to the up-Higgs in MSSM. There is a wide region of parameter space, where electroweak symmetry breaking occurs but left squarks remain heavy. In this model the favorable hierarchy

$$m_Q \gg m_{\tilde{t}_R}, A_t$$

appears naturally. The reason for the hierarchy between  $m_Q$  and  $A_t$  is that the largest contribution to the low energy value of  $A_t$  comes from two-loop diagrams involving gluino; in this model gluino is light, that reduces  $A_t$  significantly. As a consequence, a fairly wide region of  $\tan\beta$  will be consistent with the “light stop” solution in this model.

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<sup>4</sup>The same observations apply to models where messengers carrying different quantum numbers are characterized by different scales  $\Lambda$ .

## 5 Conclusions

We have demonstrated that electroweak baryogenesis may produce acceptable amount of baryon asymmetry in the framework of GMSB models with messenger-matter mixing. The required MSSM spectrum is similar to the so-called “light stop window” of SUGRA models [7]. We have plainly mapped that “window” onto the space of parameters of GMSB models. We have found that severe constraints are to be imposed on the GMSB parameters. At least one of the mixing terms has to be quite large in order to make right stop lighter than top. In the explicit example presented in this letter, the minimal possible value of CP-phase is reached at a large number of anti-symmetric messenger generations,  $n_{10} \gtrsim 24$ . With maximal CP-violation ( $\phi_{CP} \sim 1$ ) one has  $n_{10} \gtrsim 6$ . Another property of GMSB models with electroweak baryogenesis is  $\tan\beta \lesssim 4$ . We have briefly outlined also extensions of this model which have better properties with respect to electroweak baryogenesis. Let us note finally, that additional sources of CP-violation, which may originate from messenger-matter mixing, may enhance the electroweak production of baryon asymmetry.

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